Superfluid Flow of Polaron Polaritons above Landau's Critical Velocity

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We develop a theory for the interaction of light with superfluid optical media, describing the motion of quantum impurities that are created and dragged through the liquid by propagating photons. It is well known that a mobile impurity suffers dissipation due to phonon emission as soon as it moves faster than the speed of sound in the superfluid—Landau's critical velocity. Surprisingly we find that in the present hybrid light-matter setting, polaritonic impurities can be protected against environmental decoherence and be allowed to propagate well above the Landau velocity without jeopardizing the superfluid response of the medium.

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When an object moves through a superfluid it can do so without friction as long as it is slower than a certain critical velocity. In his seminal work [1], Landau obtained this bound by arguing that a moving impurity can generate excitations only when it exceeds the speed of sound in the superfluid. In this case, the object emits Cherenkov radiation which decelerates its motion. Being a hallmark of superfluidity this effect and the associated Landau velocity have since been investigated in diverse systems, from liquid helium [2–4] and exciton-polariton fluids in semiconductor microcavities [5], to ultracold atomic quantum gases [6].

An atomic impurity inside an ultracold gas of bosonic atoms [7–15] provides an ideally suited and well controllable platform to study such behavior, as demonstrated in recent experiments [16–19]. These measurements revealed the emergence of a polaron quasiparticle in close analogy to its solid-state counterpart, introduced more than 80 years ago [20,21] to understand how electrons interact with lattice vibrations of the surrounding crystal. The underlying Fröhlich model [21] has since found applications to various problems. For example, light-matter interactions originate from the optical generation of excitations in the material, whereby the coupling [21] between such excitons and phonons can lead to dissipation and explains some important optical properties of semiconductors [22]. The realization of strong light-matter coupling in such systems has enabled broad explorations of collective phenomena [5,23–28] and future applications [29–36] of exciton polaritons. However, their coupling to phonons and ensuing damping of polarons remains a major limiting factor for coherence and quantum effects in such systems.

Here, we address this issue by developing a theory for the nonequilibrium dynamics of polaritons in a quantum many-body system under the formation of Fröhlich polarons [see Fig. 1(a)]. Considering the three-level scheme illustrated in Fig. 1(b), we demonstrate the emergence of polaron-polariton quasiparticles that can vastly exceed the traditional Landau critical velocity of the medium without suffering phonon-induced decoherence [see Fig. 1(c)]. This effect, in turn, permits us to stabilize and protect an otherwise decaying polaron against phonon-induced decoherence via a vanishingly small photon component of the formed polariton [see Fig. 1(d)]. The discovery of such unusual behavior sheds new light on the optical properties of quantum many-body systems and may open up new routes for controlling and mitigating phonon-induced decoherence in light-matter interfaces.

More specifically, we consider a superfluid medium consisting of a weakly interacting atomic Bose-Einstein condensate (BEC), whereby an incident photon may transfer an atom to a different internal quantum state, which then acts as an impurity. Its interaction with the surrounding superfluid generates phonons, which screen the impurity to form a polaronic quasiparticle. To avoid dissipation from radiative decay of the excited state $|e\rangle$, one can apply an additional control field and couple two stable atomic states, the state $|b\rangle$ comprising the BEC and the state $|c\rangle$ being the impurity state, via a two-photon transition as shown in Fig. 1(b). On two-photon resonance, the depicted three-level scheme realizes electromagnetically induced transparency (EIT), which affords strong light-matter coupling at virtually vanishing photon losses [37] due to the formation of so-called dark-state polaritons [38] that propagate with a greatly reduced group velocity, v_g, as low as a few m/s [39]. At such low group velocities, the dark-state polariton is primarily composed of the impurity excitation with a very low photon fraction less than 10^{-6} [38].

Taken separately, these scenarios thus yield two stable quasiparticles: a photon-dressed impurity and a phonon-dressed impurity, which remains stable as long as its velocity

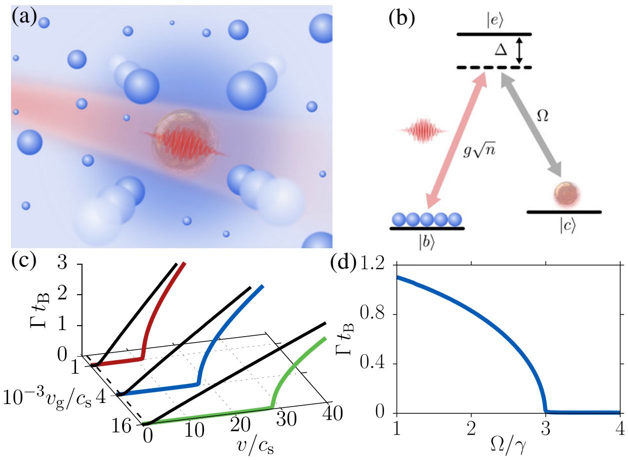

FIG. 1. (a) Illustration of a propagating photon generating a dark-state polaron polariton via impurity interactions. (b) The incident photon can form a dark-state polariton by coupling the atomic ground state $|b\rangle$ to an excited state $|e\rangle$ with a detuning Δ and a coupling strength $g\sqrt{n}$, determined by the atomic density n. The state $|e\rangle$ decays radiatively with rate γ and is coupled to a stable impurity state $|c\rangle$ via a classical control field with Rabi frequency Ω. Panels (c) and (d) show the decay rate Γ of the formed polaron polariton in units of $t_B = \xi/\sqrt{2}c_s$, determined by the coherence length, ξ, and the speed of sound, c_s, of the superfluid. (c) Γ as a function of the impurity speed v and the polariton group velocity $v_{\rm g}$, varied through the density, $n \simeq 2\times 10^{14}$ (red), 0.8×10^{14} (blue), 0.3×10^{14} cm^{-3} (green), for $\Omega/\gamma = 2$ and $\Delta/\gamma = -200$. The damping of the bare polaron is shown by the black lines. (d) Γ as a function of Ω for $n \simeq 0.8\times 10^{14}$ cm^{-3} and $\Delta/\Omega = -300$, revealing the emergence of a critical field $\Omega/\gamma \sim 3$. All calculations are performed for the D_1 transition of ultracold ^{23}Na atoms and an impurity scattering length of $a = 500a_0$, in units of the Bohr radius a_0.

is below the Landau velocity, i.e., the speed of sound in the superfluid. Consequently, one would expect that the combined quasiparticle destroys superfluidity [40] as soon as $v_{\rm g}$ exceeds Landau's critical velocity. Surprisingly, this is not the case. First, it turns out that it is not the group velocity which determines the viscosity of its environment, but the total recoil momentum exerted on the impurity state by the two applied light fields. The resulting impurity velocity v, is widely tunable via the angle between the two laser fields and can differ vastly from $v_{\rm g}$. Second, we show that *both* of these velocities of the moving impurity can greatly exceed Landau's critical velocity without destroying the superfluid response of the quantum liquid [see Fig. 1(c)].

In order to understand these findings, let us consider a BEC of atoms with a mass m, a density n, and three internal states $|b\rangle$, $|e\rangle$, and $|c\rangle$, which are coupled by the propagating quantum light field and a classical control laser as indicated in Fig. 1(b). We focus on weak collisional interactions that are short ranged and can be parametrized by a scattering length a_B for the condensate atoms in the ground state $|b\rangle$ and a scattering length a quantifying the interaction between the impurity atoms in the $|c\rangle$ state and the condensate. The underlying Hamiltonian $\hat{H} = \hat{H}_0 + \hat{H}_{\rm int} + \hat{H}_{\rm al}$ [41] can be conveniently split into three parts. Here,

$$\hat{H}_0 = \sum_{\mathbf{p}} \varepsilon^{\alpha}_{\mathbf{p}} \hat{\alpha}^{\dagger}_{\mathbf{p}} \hat{\alpha}_{\mathbf{p}} + \sum_{\mathbf{k}} [\varepsilon^{e}_{\mathbf{k}} \hat{e}^{\dagger}_{\mathbf{k}} \hat{e}_{\mathbf{k}} + \varepsilon^{c}_{\mathbf{k}} \hat{c}^{\dagger}_{\mathbf{k}} \hat{c}_{\mathbf{k}} + \omega_{\mathbf{k}} \hat{\beta}^{\dagger}_{\mathbf{k}} \hat{\beta}_{\mathbf{k}}] \quad (1)$$

describes one-body energies of the incident photons and the atoms in the atomic states $|e\rangle$, and $|c\rangle$, which are respectively created by the operators $\hat{\alpha}^{\dagger}_{\mathbf{p}}$ for a given momentum $\mathbf{p}$, and $\hat{e}^{\dagger}_{\mathbf{k}}$, $\hat{c}^{\dagger}_{\mathbf{k}}$ with a given momentum $\mathbf{k}$. We consider a narrow-band incoming photon field, propagating along the z axis with momenta $\mathbf{p}$ that are tightly centered around the carrier momentum $\mathbf{p}_0 = p_0 \mathbf{e}_z$. This defines a rotating frame in which the photon energy is $\varepsilon^{\alpha}_{\mathbf{p}} = c(p - p_0)$, with the speed of light c. The complex energy $\varepsilon^{e}_{\mathbf{k}} = k^2/2m + \Delta - i\gamma$ of excited-state atoms contains the one-photon detuning Δ and decay rate γ, while the energy $\varepsilon^{c}_{\mathbf{k}} = k^2/2m + \delta$ of the impurity state is set by the two-photon detuning δ. Excitations of the weakly interacting condensate are Bogoliubov modes, created by $\hat{\beta}^{\dagger}_{\mathbf{k}} = u_{\mathbf{k}} \hat{b}^{\dagger}_{\mathbf{k}} + v_{\mathbf{k}} \hat{b}_{-\mathbf{k}}$ at momenta $\mathbf{k}$ with energy $\omega_{\mathbf{k}}$, whereby $b^{\dagger}_{\mathbf{k}}$ creates an atom in the atomic ground state $|b\rangle$ and $u_{\mathbf{k}}$, $v_{\mathbf{k}}$ are the corresponding BEC coherence factors [42]. The light-matter interaction,

$$\hat{H}_{\rm al} = \Omega \sum_{\mathbf{k}} \hat{c}^{\dagger}_{\mathbf{k}-\mathbf{k}_{\rm cl}} \hat{e}_{\mathbf{k}} + \frac{g}{\sqrt{V}} \sum_{\mathbf{k},\mathbf{p}} \hat{b}^{\dagger}_{\mathbf{k}} \hat{\alpha}^{\dagger}_{\mathbf{p}} \hat{e}_{\mathbf{p}+\mathbf{k}} + \text{H.c.}, \quad (2)$$

describes the coupling to the classical control field with wave vector $\mathbf{k}_{\rm cl}$ and Rabi frequency Ω, as well as the single-photon interaction with a coupling strength g within the rotating wave approximation. While the sum over $\mathbf{p}$ is restricted to momenta for the incident photons, the photonic vacuum has been integrated out [41] yielding the decay rate γ of the excited state included in $\varepsilon^{e}_{\mathbf{k}}$ above. In the absence of atomic interactions and at the two-photon resonance $\delta = 0$, the dynamics governed by $\hat{H}_0 + \hat{H}_{\rm al}$ shows that incoming photons are converted to dark-state polaritons $\hat{d}_{\mathbf{p}} = \cos\theta \hat{\alpha}_{\mathbf{p}} - \sin\theta \hat{c}_{\mathbf{p}-\mathbf{k}_{\rm cl}}$ that propagate the medium without losses at a velocity $v_{\rm g} = \cos^2\theta c$, determined by $\tan\theta = g\sqrt{n}/\Omega$ [38]. The typical case of large single-photon Rabi frequencies $g\sqrt{n} \gg \Omega$ [39], thus effectively yields an impurity $\hat{c}_{\mathbf{p}-\mathbf{k}_{\rm cl}} \approx -\hat{d}_{\mathbf{p}}$ that has a form stable propagation through the condensate with an ultraslow velocity $v_{\rm g} \ll c$.

The interaction between the impurity and the superfluid can be described by the Fröhlich Hamiltonian [21]

$$\hat{H}_{\rm int} = \frac{\sqrt{n}\mathcal{T}}{\sqrt{V}} \sum_{\mathbf{q},\mathbf{k}} (u_{\mathbf{k}} - v_{\mathbf{k}}) \hat{c}^{\dagger}_{\mathbf{q}-\mathbf{k}} \hat{c}_{\mathbf{q}} (\hat{\beta}^{\dagger}_{\mathbf{k}} + \hat{\beta}_{-\mathbf{k}}), \quad (3)$$

which serves as a paradigmatic model for a range of solid-state systems [21,43] and applies to polarons in

BECs with weak interactions [13]. Physically, Eq. (3) describes momentum-changing impurity collisions that generate Bogoliubov excitations with an underlying scattering matrix $\mathcal{T} = 4\pi a/m$. These collisions can profoundly alter the idealized scenario of dissipation-free polariton motion.

To characterize the resulting many-body dynamics, we use an ansatz

$$|\Psi_{\mathbf{p}}(t)\rangle = [A_{\mathbf{p}}^{(0)}(t)\hat{\alpha}_{\mathbf{p}}^{\dagger} + E_{\mathbf{p}}^{(0)}(t)\hat{e}_{\mathbf{p}}^{\dagger} + C_{\mathbf{p}}^{(0)}(t)\hat{c}_{\mathbf{p}-\mathbf{k}_{\mathrm{cl}}}^{\dagger}]|\mathrm{BEC}\rangle + \sum_{\mathbf{k}}[E_{\mathbf{p},\mathbf{k}}^{(1)}(t)\hat{e}_{\mathbf{p}-\mathbf{k}}^{\dagger} + C_{\mathbf{p},\mathbf{k}}^{(1)}(t)\hat{c}_{\mathbf{p}-\mathbf{k}_{\mathrm{cl}}-\mathbf{k}}^{\dagger}]\hat{\beta}_{\mathbf{k}}^{\dagger}|\mathrm{BEC}\rangle, \quad (4)$$

for the time-dependent wave function, which is truncated at the single phonon level to leading order in the impurity interaction. Here $|\mathrm{BEC}\rangle$ denotes the initial state of the Bose-Einstein condensate composed entirely of $|b\rangle$-state atoms. The first line describes the bare photon-driven impurity dynamics that yields the lossless propagation of the dark-state polariton amplitude $D_{\mathbf{p}} = \langle \mathrm{BEC}|d_{\mathbf{p}}|\Psi_{\mathbf{p}}(t)\rangle = \cos\theta A_{\mathbf{p}}^{(0)} - \sin\theta C_{\mathbf{p}}^{(0)}$ discussed above. Collisions between the impurity and the surrounding atoms, however, perturb this polariton state and excite the superfluid as described by the Fröhlich term in Eq. (3) and captured by the second line in Eq. (4). The characteristic momentum change associated with such collisions is given by the inverse coherence length $1/\xi = \sqrt{8\pi n a_B}$ of the condensate, which for a large single-photon detuning, $|\Delta| \gg \gamma$, lies far outside the EIT regime. Consequently, almost all impurity collisions, apart from negligible scattering events around $|\mathbf{p}-\mathbf{k}| \simeq p$ [41], lead to a break up of the low-energy dark-state polariton and populate the hybridized modes $|\pm\rangle$ of the two laser-coupled $|e\rangle$ and $|c\rangle$ states with energies $\varepsilon_{\mathbf{p}}^{(\pm)} = [\varepsilon_{\mathbf{p}}^{e} + \varepsilon_{\mathbf{p}-\mathbf{k}\mathrm{cl}}^{c} \pm (4\Omega^2 + (\varepsilon_{\mathbf{p}}^{e} - \varepsilon_{\mathbf{p}-\mathbf{k}\mathrm{cl}}^{c})^2)^{1/2}]/2$ as indicated in Figs. 2(a) and 2(b). This implies a prompt photon loss and is reflected in the omission of the photon component in the second line of Eq. (4). It is this interaction-induced modification of the polariton character and associated dispersions that causes the unusual propagation phenomena found in this Letter.

By using this ansatz in the many-body Schrödinger equation $i\partial_t|\Psi_{\mathbf{p}}\rangle = \hat{H}|\Psi_{\mathbf{p}}\rangle$ we obtain a set of coupled equations for the five state amplitudes in Eq. (4). Upon solving the evolution equations for $E^{(1)}$ and $C^{(1)}$ and substituting the result into the equations for the zero phonon amplitudes, we derive a closed equation [41]

$$i\partial_t D_{\mathbf{p}}(t) = [\varepsilon_{\mathbf{p}} + \Sigma_{\mathbf{p}} - \tilde{\Sigma}_{\mathbf{p}}(t)]D_{\mathbf{p}}(t) \quad (5)$$

that describes the open quantum dynamics of the dark-state polariton due to its interaction with the surrounding superfluid. Here, $\varepsilon_{\mathbf{p}} = v_{\mathrm{g}}(p - p_0) + \sin^2\theta(\mathbf{p} - \mathbf{k}_{\mathrm{cl}})^2/2m$ is the dispersion of the noninteracting dark state polariton around

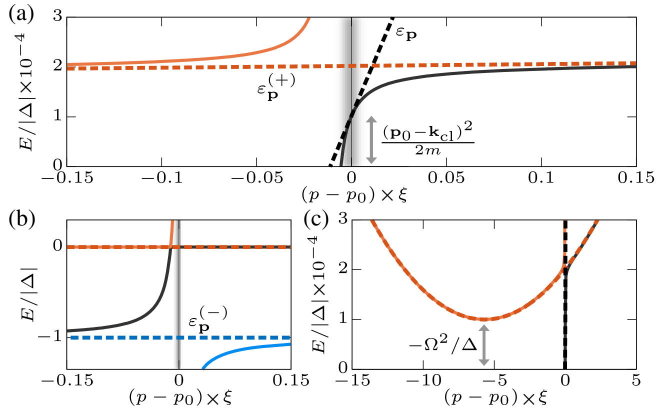

FIG. 2. Polariton dispersion curves in the absence of atomic interactions for $\Delta = -200\gamma$, $\Omega = 2\gamma$, and $n = 0.5\times10^{14}\ \mathrm{cm}^{-3}$. (a) Incoming photons generate dark-state polaritons (black solid line) with an approximate linear dispersion, $\varepsilon_{\mathbf{p}} \simeq v_{\mathrm{g}}(p-p_0) + (\mathbf{p}-\mathbf{k}_{\mathrm{cl}})^2/2m$, around $p - p_0 \approx 0$ (black dotted line) that facilitates low-loss form stable photon propagation with the slow-light group velocity v_{g}. Atomic collisions with the surrounding condensate cause a typical momentum change of $\Delta p \sim 1/\xi$ well outside this EIT regime, indicated by the vertical grey bar. The dark state is thereby broken apart by any atomic collision event, and scatters into the photon-free hybridized states $|\pm\rangle$, with the indicated energies $\varepsilon_{\mathbf{p}}^{(\pm)}$, shown by the orange and blue dashed lines in panels (a) and (b). This characteristic scattering process leads to the ansatz Eq. (4) for the polaron polariton. As illustrated in panel (b), the energy of the state $|-\rangle$ is typically so far removed that it does not contribute significantly to the emerging polaron-polariton quasiparticle and its self-energy, Eq. (6). Panel (c) shows the same dispersion curves on an expanded momentum scale, revealing the quadratic contribution from the atomic kinetic energy and the light shift induced by the classical control field.

p_0 [see Fig. 2(a)]. The second term accounts for the kinetic energy of the atoms and is normally discarded when describing slow-light propagation [37,38]. Here, however, it plays a crucial role in capturing the physics of atomic interactions. The time-dependent complex energy $\tilde{\Sigma}_{\mathbf{p}}(t)$ [41] captures the nonequilibrium dynamics driven by the atomic interactions following the creation of the ideal dark state polariton at time $t = 0$. The vanishing of $\tilde{\Sigma}_{\mathbf{p}}(t)$ at longer times then signals the establishment of a new quasiparticle—the polaron polariton. Its self-energy

$$\Sigma_{\mathbf{p}} = \int \frac{d^3k}{(2\pi)^3}\left[\frac{(g_{\mathbf{p},\mathbf{k}}^{(+)})^2}{\varepsilon_{\mathbf{p}} - \varepsilon_{\mathbf{p}-\mathbf{k}}^{(+)} - \omega_{\mathbf{k}}} + \frac{(g_{\mathbf{p},\mathbf{k}}^{(-)})^2}{\varepsilon_{\mathbf{p}} - \varepsilon_{\mathbf{p}-\mathbf{k}}^{(-)} - \omega_{\mathbf{k}}} + \sin^2\theta\cdot n\mathcal{T}^2\frac{m}{k^2}\right] \quad (6)$$

describes the effects of interactions on the quasiparticle dispersion and has a simple physical interpretation. First note that the classical control field hybridizes the $|e\rangle$ and $|c\rangle$ states of the atoms and generates new dressed states $|\pm\rangle$

with energies $\varepsilon_{\mathbf{p}}^{(\pm)}$, as outlined above and indicated in Fig. 2. Equation (6) therefore describes the virtual scattering of the impurity into these hybridized modes $|\pm\rangle$ upon the generation of phonon excitations with an energy $\omega_{\mathbf{k}}$. The associated coupling elements [41]

$$g_{\mathbf{p},\mathbf{k}}^{(+)} = \sin\theta\left[u_{\mathbf{p}-\mathbf{k}}^{\mathrm{ec}}\sqrt{n}\mathcal{T}(u_{\mathbf{k}} - v_{\mathbf{k}}) + w_{\mathbf{p}-\mathbf{k}}^{\mathrm{ec}}\frac{v_{\mathbf{k}}\Omega}{\sqrt{n}}\right],$$
$$g_{\mathbf{p},\mathbf{k}}^{(-)} = \sin\theta\left[w_{\mathbf{p}-\mathbf{k}}^{\mathrm{ec}}\sqrt{n}\mathcal{T}(u_{\mathbf{k}} - v_{\mathbf{k}}) - u_{\mathbf{p}-\mathbf{k}}^{\mathrm{ec}}\frac{v_{\mathbf{k}}\Omega}{\sqrt{n}}\right] \tag{7}$$

are determined by the form of the hybridized states, described by $u_{\mathbf{q}}^{\mathrm{ec}} = (\varepsilon_{\mathbf{q}}^{(+)} - \varepsilon_{\mathbf{q}}^{e})/[(\varepsilon_{\mathbf{q}}^{(+)} - \varepsilon_{\mathbf{q}}^{e})^2 + \Omega^2]^{1/2}$ and $w_{\mathbf{q}}^{\mathrm{ec}} = \Omega/[(\varepsilon_{\mathbf{q}}^{(+)} - \varepsilon_{\mathbf{q}}^{e})^2 + \Omega^2]^{1/2}$, whereby $g_{\mathbf{p},k}^{(-)}$ vanishes as $g_{\mathbf{p},k}^{(-)} \sim \Omega$ with a decreasing control field. Eventually, Eq. (6) approaches the known second order polaron energy [9] in the zero-field limit in which the dark-state polariton coincides with the bare impurity. The obtained equation of motion, Eq. (5), has a simple solution $D_{\mathbf{p}}(t) = D_{\mathbf{p}}(0)\mathrm{e}^{-iE_{\mathbf{p}}t - \Gamma_{\mathbf{p}}t}\mathrm{e}^{i\int_0^t d\tau\tilde{\Sigma}_{\mathbf{p}}(\tau)}$. Starting from an initially noninteracting dark-state polariton, $D_{\mathbf{p}}(0)$, this solution describes the initial quasiparticle formation, as determined by $\tilde{\Sigma}_{\mathbf{p}}(t)$, and the subsequent evolution of the formed polaron polariton, governed by its energy $E_{\mathbf{p}} = \varepsilon_{\mathbf{p}} + \mathrm{Re}\Sigma_{\mathbf{p}}$ and steady-state damping rate $\Gamma_{\mathbf{p}} = -\mathrm{Im}\Sigma_{\mathbf{p}}$. In the more familiar case of a bare polaron ($\Omega = 0$), the impurity suffers a finite damping rate, $\Gamma_{\mathbf{p}}$, if it moves faster than the Landau critical velocity, given by the condensate's speed of sound $c_s = \sqrt{4\pi a_B n}/m$. The kinetic energy is then sufficient to generate phonon excitations with a low-energy dispersion $\omega_{\mathbf{k}} \simeq c_s k$ and cause dissipation in the form of Cherenkov radiation [1]. However, the damping rate of our dark-state polaron polariton, shown in Fig. 1(c), suggests profoundly different behavior than this paradigmatic scenario for the breakdown of superfluidity.

We observe that the group velocity, v_g, which governs the speed with which the impurity excitation traverses the medium, has virtually no bearing on the damping of the polaron and can exceed c_s by several orders of magnitude. In fact, it turns out that it is not the velocity v_g of the polaritonic quasiparticle that determines the superfluid response of the medium, but the velocity of the laser-excited impurity atom. This velocity, $\mathbf{v} = (\mathbf{p} - \mathbf{k}_{\mathrm{cl}})/m$, can be widely tuned via the propagation angle between the incident control laser and the probe photons with wave vectors $\mathbf{k}_{\mathrm{cl}}$ and $\mathbf{p} \simeq \mathbf{p}_0$, respectively.

Yet, even this velocity can exceed the speed of sound of the condensate by more than an order of magnitude without jeopardizing its superfluid response, as shown in Fig. 1(c). To understand this behavior, we consider the off-resonant limit, $\Omega/|\Delta| \ll 1$, in which the $|-\rangle$ state is far removed in energy as shown in Fig. 2(b), whereby the term involving $g_{\mathbf{p},\mathbf{k}}^{(-)}$ in Eq. (6) can be neglected. As a result, the denominator of the first term in Eq. (6) dictates the energy balance

$$\frac{(\mathbf{p} - \mathbf{k}_{\mathrm{cl}})^2}{2m} = \frac{(\mathbf{p} - \mathbf{k}_{\mathrm{cl}} - \mathbf{k})^2}{2m} + \omega_{\mathbf{k}} - \frac{\Omega^2}{\Delta}, \tag{8}$$

for the scattering of a polariton with energy $\varepsilon_{\mathbf{p}}$ into a different momentum state with $\varepsilon_{\mathbf{p}-\mathbf{k}}^{(+)} \simeq (\mathbf{p} - \mathbf{k}_{\mathrm{cl}} - \mathbf{k})^2/2m - \Omega^2/\Delta$ while emitting a phonon with an energy $\omega_{\mathbf{k}}$ via collisions between the impurity and its surrounding atoms. To obtain Eq. (8), we set $\varepsilon_{\mathbf{p}} \simeq (\mathbf{p} - \mathbf{k}_{\mathrm{cl}})^2/2m$, since $\sin\theta \simeq 1$, and because the photon momentum p is well within the EIT window such that $v_g|p - p_0| \ll \Omega^2/|\Delta|$ [see Fig. 2(a)]. Without the light field ($\Omega = 0$), Eq. (8) permits phonon emission only for impurity velocities $v = |\mathbf{p} - \mathbf{k}_{\mathrm{cl}}|/m \geq c_s$ above the familiar Landau critical velocity $v_c = c_s$. In contrast, the presence of the light field renders the impurity collisions inelastic by introducing an additional energy cost $-\Omega^2/\Delta$ associated with the collisional break up of the dark-state polariton into the laser-dressed $|+\rangle$-state impurity as indicated in Fig. 2(c). For a positive single-photon detuning, $\Delta > 0$, the resulting endothermic character of the impurity collisions promotes phonon emission regardless of the impurity speed, corresponding to a vanishing critical velocity, $v_c = 0$.

A negative detuning, $\Delta < 0$, on the other hand, introduces an additional energy cost for impurity collisions and thereby increases the critical velocity. Upon increasing the light shift Ω^2/Δ, this effect can indeed cause a substantial enhancement and increase the critical velocity by more than an order of magnitude under typical conditions of ultracold atom experiments [39]. At the same time, this effect enables the quantum optical stabilization of otherwise decaying polaron quasiparticles. Indeed, Fig. 1(d) reveals the emergence of a critical behavior with respect to the control field amplitude and demonstrates the efficient protection of the polaron against the otherwise inevitable emission of Cherenkov radiation above a critical control field $\Omega_c \simeq v\sqrt{-m\Delta/4}$ [41].

This optical stabilization of the Bose polaron against phonon emission can be probed directly by measuring the transmission of slow-light polaritons through an ultracold gas of Bose condensed atoms. The propagation dynamics through the gas is conveniently visualized by Fourier transforming the obtained solution, $D_{\mathbf{p}}(t)$, into real space. Figure 3 compares the resulting pulse evolution for a bare Bose polaron and a dark-state polaron polariton, moving at initially identical velocities through a ^{23}Na condensate with experimentally accessible densities and laser parameters. The Bose polaron undergoes rapid decoherence due to the steady emission of Cherenkov radiation [41], while the amplitude of the dark-state polaron polariton settles at the quasiparticle residue [44] and remains otherwise protected

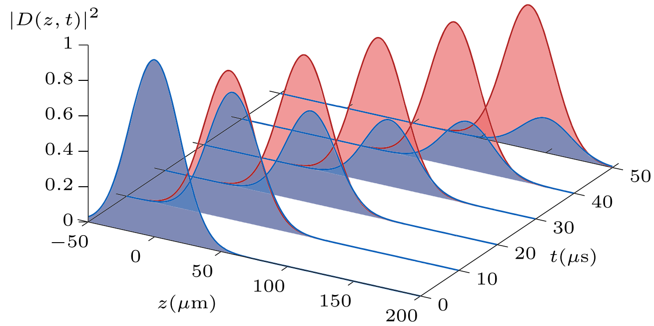

FIG. 3. Pulse propagation through a condensate of ^{23}Na atoms with a density of $n = 2.6 \times 10^{14}$ cm^{-3} and an impurity scattering length of $a = 0.1\xi$. The dynamics of a bare impurity wave packet (blue lines) suffers strong damping due to the supersonic motion of the formed polaron with an initial velocity of 3 m/s $\gg c_s$. In contrast, the red lines show the asymptotically undamped motion of a polaron polariton with $\Delta = -200\gamma$ and an identical initial group velocity of $v_g = 3$ m/s, corresponding to a near-unity impurity fraction of $1 - v_g/c = 0.99999999$. Polaron formation eventually leads to a slight lowering of the group velocity [41].

from decoherence, eventually propagating at a lowered group velocity $v_g + \partial_p \mathrm{Re}\Sigma_{\mathbf{p}}|_{\mathbf{p}_0}$.

The demonstrated ability to stabilize mobile polaritons in a dissipative environment thus provides an intriguing outlook for realizing coherent optical interfaces and makes it possible to explore and control the combined formation of polaritonic and polaronic quasiparticle states at greatly reduced losses and decoherence. Not only does this combination yield an attractive platform for exploring impurity physics [40], and suggest novel optical probes of quantum many-body dynamics [45], but also promises new functionalities for light-matter interfaces and optical devices [46,47]. In the present context, ensuing applications include the generation of few-photon nonlinearities via induced polaron interactions in atomic superfluids [48,49], which may even be controlled and enhanced via resonant phonon-exchange processes. Moreover, as outlined above, the underlying interaction Hamiltonian [Eq. (3)] is of considerably greater applicability describing for example the coupling between excitons and phonons in semiconductors [43], which often presents a limitation to the coherence of light-matter interactions in such systems [22]. The EIT-enabled stabilization against phonon-induced dissipation, described in this Letter, therefore suggests a promising approach to alleviating this obstacle. These combined perspectives motivate future investigations into the strong-coupling regime as well as a wider range of environmental interactions and photon interfaces for exploiting correlated quantum dynamics and exploring quantum nonlinear optics in strongly interacting many-body systems.

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